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Anomalous $\gamma\gamma H$ and $Z\gamma H$ couplings in the process $e\gamma \to eH$

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Abstract

The bounds on the anomalous contributions to the $\gamma\gamma H$ and $Z\gamma H$ vertices that can be obtained via the process $e\gamma \to eH$ are discussed through a model independent analysis based on $SU(2)\times U(1)$ invariant operators of $dim = 6$. A light Higgs boson with $m_H \simeq 120$ GeV and center-of-mass energies of $\sqrt{S} = 500$ GeV and $\sqrt{S} = 1500$ GeV are assumed. The process $e\gamma \to eH$ is shown to provide an excellent way to strongly constraint both the CP-even and CP-odd anomalous contributions in the $\gamma\gamma H$ and $Z\gamma H$ couplings.

1 Introduction

Once the Higgs boson will be discovered, the next linear e^+e^- colliders with centre-of-mass (c.m.) energies $\sqrt{s} \simeq (300 \div 2000)$GeV and integrated luminosity $\mathcal{O}(100-1000)$ fb^{-1} will allow an accurate determination of its mass, some couplings and parity properties [1, 2]. Among other couplings, the interaction of the Higgs scalar with the neutral electroweak gauge bosons, γ and Z, are particularly interesting. Indeed, an accurate determination of these couplings will allow to test some delicate feature of the Standard Model: the relation between the spontaneous symmetry breaking mechanism and the electroweak mixing of the two gauge groups $SU(2)$ and $U(1)$. In this respect, three vertices could be measured, ZZH, $\gamma\gamma H$ and $Z\gamma H$. In the SM the ZZH vertex stands at the tree level, while the other two are generated at one-loop. Therefore these last ones are particularly interesting, also

because the effective $\gamma\gamma H$ and $Z\gamma H$ couplings could be sensitive to the contributions of new heavy particles circulating in the loop.

Here, we consider the case of a light Higgs boson, that is with $M_Z \lesssim m_H \lesssim 140$ GeV. For this range of mass, a measurement of the $\gamma\gamma H$ coupling should be possible by the determination of the BR for the decay $H \to \gamma\gamma$ in the channel $gg \to H \to \gamma\gamma$ at the next large hadron collider LHC (see, for instance, [3]). Furthermore, at future photon-photon colliders*, the precise measurement of the $\gamma\gamma H$ vertex can be achieved at the resonant production of the Higgs particle, $\gamma\gamma \to H$. To this end, the capability of tuning the $\gamma\gamma$ c.m. energy on the Higgs mass, through a good degree of the photons monochromaticity, will be crucial for not diluting too much the $\gamma\gamma \to H$ resonant cross section over the c.m. energy spectrum.

Measuring the $Z\gamma H$ vertex is in general a more complicated task. One possibility to test this vertex is given by the $H \to \gamma Z$ decay. Unfortunately, this decay suffers from having a small rate, and moreover, for the main Z decays, the signal is affected by large backgrounds.

Another possibility of measuring the $Z\gamma H$ vertex is given by Higgs production processes. At electron-positron colliders, the corresponding channels are $e^+e^- \to \gamma H$ and $e^+e^- \to ZH$. The reaction $e^+e^- \to \gamma H$ has been extensively studied in the literature [6, 7, 8]. Unfortunately, the $e^+e^- \to H\gamma$ channel suffers from small rates, which are further depleted at large energies by the $1/s$ behavior of the dominant s-channel diagrams. The crossed process, the one-loop Higgs production in electron-photon collisions through $e\gamma \to eH$, was analysed in [9, 10, 11]. This channel turns out to be a golden place to test both the $\gamma\gamma H$ and $Z\gamma H$ one-loop couplings with high statistics, without requiring a fine tuning of the c.m. energy. Indeed, the total rate of this reaction is rather high [9]. In particular, for m_H up to about 400 GeV, one finds $\sigma > 1$ fb. If no kinematical cuts are imposed, then the main contribution to the cross section is given by the $\gamma\gamma H$ vertex. On the contrary, the $Z\gamma H$ vertex contribution is depleted by the Z propagator. Nevertheless, the $Z\gamma H$ vertex effects can be extracted from $e\gamma \to eH$ by implementing a suitable strategy to reduce the $\gamma\gamma H$ vertex contribution [9]. This strategy requires a final electron tagged at large angle together with a transverse momentum cut $p_T^e > 100$ GeV. The main irreducible background to the process $e\gamma \to eH \to eb\bar{b}$ comes from the channel $e\gamma \to eb\bar{b}$. However, proper cuts on the quark-pair invariant mass and on the angles $\theta_{b(c)-beam}$ (between each b quark direction and both the beams) can bring the signal and background at a comparable level. Then, with a luminosity of 100 fb^{-1}, at $\sqrt{s} = 500$GeV, one expects an accuracy as good as about 10% on the measurement of the $Z\gamma H$ effects assuming the validity of the SM [9]. †

*Two further options are presently considered for a high-energy e^+e^- linear collider, where one or both the initial e^+/e^- beams are replaced by photon beams induced by Compton backscattering of laser light on the high-energy electron beams [4]. Then, the initial real photons could be to a good degree monochromatic, and have energy and luminosity comparable to the ones of the parent electron beam [5].

†Since we are discussing the measurement of $\gamma\gamma H$ and $Z\gamma H$ couplings at the $\mathcal{O}(10\%)$ level, the QCD

Here, we present the prospects of the $e\gamma \to eH$ reaction in setting experimental bounds on the value of possible anomalous $\gamma\gamma H$ and $Z\gamma H$ couplings [13]. Some preliminary results have been presented in [14], too (see also [15]). For this analysis we use a model independent approach, where $dim = 6$ $SU(2) \times U(1)$ invariant operators are added to the SM Lagrangian. In realistic models extending the SM, these operators contribute in some definite combinations. However, if one discusses the bounds on possible deviations from the SM one-loop Higgs vertices, this approach can give some general insight into the problem. These anomalous operators contribute to the three vertices $\gamma\gamma H$, $Z\gamma H$ and ZZH, with only the first two involved in the $e\gamma \to eH$ reaction.

We show that the reaction $e\gamma \to eH$ has an excellent potential in bounding both the vertices $\gamma\gamma H$ and $Z\gamma H$. While in the case of $\gamma\gamma H$, it is complementary to the resonant $\gamma\gamma \to H$ reaction, $e\gamma \to eH$ offers a unique possibility to investigate possible anomalies also in the $Z\gamma H$ couplings.

2 Anomalous vertices

We consider the possibility that the new physics affects the bosonic sector of the SM through low-energy effective-operators of $dim = 6$. We restrict our analysis to the set of operators that contribute to the $e\gamma \to eH$ amplitude via anomalous couplings to the $\gamma\gamma H$ and $Z\gamma H$ vertices. This set can be divided in two pairs of $dim = 6$ operators[‡], CP-even and CP-odd respectively, giving anomalous contributions to the process $e\gamma \to eH$. In terms of these operators the most general lagrangian is given by

$$\mathcal{L}^{eff} = \mathrm{d} \cdot \mathcal{O}_{UW} + \mathrm{d_B} \cdot \mathcal{O}_{UB} + \overline{\mathrm{d}} \cdot \bar{\mathcal{O}}_{UW} + \overline{\mathrm{d}}_{\mathrm{B}} \cdot \bar{\mathcal{O}}_{UB}, \tag{1}$$

$$\mathcal{O}_{UW} = \frac{1}{v^2}\left(|\Phi|^2 - \frac{v^2}{2}\right) \cdot W^{i\mu\nu} W^i_{\mu\nu}, \qquad \mathcal{O}_{UB} = \frac{1}{v^2}\left(|\Phi|^2 - \frac{v^2}{2}\right) \cdot B^{\mu\nu} B_{\mu\nu}, \tag{2}$$

$$\bar{\mathcal{O}}_{UW} = \frac{1}{v^2}|\Phi|^2 \cdot W^{i\mu\nu} \tilde{W}^i_{\mu\nu}, \qquad \bar{\mathcal{O}}_{UB} = \frac{1}{v^2}|\Phi|^2 \cdot B^{\mu\nu} \tilde{B}_{\mu\nu}, \tag{3}$$

where $\tilde{W}^i_{\mu\nu} = \epsilon_{\mu\nu\mu'\nu'} \cdot W^{i\mu'\nu'}$ and $\tilde{B}_{\mu\nu} = \epsilon_{\mu\nu\mu'\nu'} \cdot B^{\mu'\nu'}$. In these formulas, Φ is the Higgs doublet and v is the electroweak vacuum expectation value. Finally, the $\gamma\gamma H$ and

corrections to the corresponding vertices could be substantial, and should be taken into account in a precise simulation of the background in future experiments. Indeed, the next-to-leading QCD corrections to the $\gamma\gamma H$ vertex in the SM were found to be positive, and at the level of a few per cent for the Higgs masses discussed here. On the other hand, one can neglect the NLO corrections to the $Z\gamma H$ vertex. One can find the corresponding detailed discussion, e.g., in [12].

[‡]We assume that the $SU(2) \times U(1)$ local gauge invariance of the SM should be valid as well as the so-called custodial symmetry of the gauge and Higgs sectors that holds in the SM [16]

$Z\gamma H$ anomalous contributions to the helicity amplitudes of $e\gamma \to eH$ are given by:

$$M_{anom}(\sigma,\lambda) = M^{\gamma\gamma}(\sigma,\lambda) + M^{\gamma Z}(\sigma,\lambda), \tag{4}$$

where

$$\begin{aligned} M^{\gamma\gamma}(\sigma,\lambda) &= \frac{4\pi\alpha}{M_Z(-t)}\sqrt{-\frac{t}{2}}\{d_{\gamma\gamma}[(u-s)-\sigma\lambda(u+s)-i\bar{d}_{\gamma\gamma}[\lambda(u-s)+\sigma(u+s)]\}, \\ M^{\gamma Z}(\sigma,\lambda) &= \frac{4\pi\alpha(-g_e^\sigma)}{M_Z(M_Z^2-t)}\sqrt{-\frac{t}{2}}\{d_{\gamma Z}[(u-s)-\sigma\lambda(u+s)-i\bar{d}_{\gamma Z}[\lambda(u-s)+\sigma(u+s)]\}. \end{aligned}$$

Here, s, t and u are the Mandelstam kinematical variables (defined as in [9]), $\sigma/2 = \pm 1/2$ and $\lambda = \pm 1$ are the electron and photon helicities, respectively. The Z charge of the electron is denoted as g_e^σ. The anomalous couplings $\mathrm{d}, \mathrm{d_B}$ (CP-even) and $\overline{\mathrm{d}}, \overline{\mathrm{d}}_\mathrm{B}$ (CP-odd) contribute to the $\gamma\gamma H$, ZZH , and $Z\gamma H$ interactions in the combinations

$$\begin{aligned} \mathrm{d}_{\gamma\gamma} &= \tan\theta_W\ \mathrm{d} + (\tan\theta_W)^{-1}\ \mathrm{d_B}, \quad \mathrm{d}_{\gamma Z} = \mathrm{d} - \mathrm{d_B} \\ \mathrm{d}_{ZZ} &= (\tan\theta_W)^{-1}\ \mathrm{d} + \tan\theta_W\ \mathrm{d_B}, \\ \overline{\mathrm{d}}_{\gamma\gamma} &= \tan\theta_W\ \overline{\mathrm{d}} + (\tan\theta_W)^{-1}\ \overline{\mathrm{d}}_\mathrm{B}, \quad \overline{\mathrm{d}}_{\gamma Z} = \overline{\mathrm{d}} - \overline{\mathrm{d}}_\mathrm{B} \\ \overline{\mathrm{d}}_{ZZ} &= (\tan\theta_W)^{-1}\ \overline{\mathrm{d}} + \tan\theta_W\ \overline{\mathrm{d}}_\mathrm{B}, \end{aligned} \tag{5}$$

where θ_W is the Weinberg angle (we assume $\sin^2\theta_W = 0.2247$) §.

3 Bounds on anomalous $\gamma\gamma H$ and $Z\gamma H$ couplings

In this section, we present the numerical results for the bounds on the anomalous couplings d, $\mathrm{d_B}$ and $\overline{\mathrm{d}}$, $\overline{\mathrm{d}}_\mathrm{B}$ which are obtained from the $e\gamma \to eH$ process. ¶ These bounds have been computed by using the requirement that no deviation from the SM cross section is observed at the 95% CL. In particular, we require [13]:

$$N^{\mathrm{anom}}(\kappa) < 1.96 \cdot \sqrt{N^{\mathrm{tot}}(\kappa)}, \quad \kappa = \mathrm{d}, \mathrm{d_B}, \overline{\mathrm{d}}\ , \overline{\mathrm{d}}_\mathrm{B}\ , \tag{6}$$

$$N^{\mathrm{tot}}(\kappa) = \mathcal{L}_{int}\cdot[\sigma_S(\kappa)+\sigma_B]\ , \quad N^{\mathrm{anom}}(\kappa) = \mathcal{L}_{int}\cdot[\sigma_S(\kappa)-\sigma_S(0)]\ . \tag{7}$$

where $\mathcal{L}_{int}$ is the integrated luminosity, N^{tot} and N^{anom} denote respectively the total number of observed events and the anomalous number of events deviating from the expected SM predictions for the signal. By $\sigma_S(\kappa)$ we mean the cross section of the signal reaction $e\gamma \to eH \to eb\bar{b}$ with the anomalous contributions. $\sigma_S(0)$ is the SM cross section. By σ_B, we denote the cross section of the background processes $e\gamma \to eb\bar{b}$, $e\gamma \to ec\bar{c}$ (with 10% probability of misidentifying a c quark into a b quark). For the kinematical cuts applied see [13].

§The adopted values for the other physical parameters of the SM are described in [9].

¶ Most of the results presented in this work were obtained with the help of the CompHEP package [17].

	$\sqrt{S} = 500$GeV	$\sqrt{S} = 500$ GeV	$\sqrt{S} = 1500$GeV	$\sqrt{S} = 1500$ GeV
$P_e = 0$	$p_T^e > 0$	$p_T^e > 100$ GeV	$p_T^e > 0$	$p_T^e > 100$ GeV
$\mathrm{d} \times 10^3$	$(-0.73,\ 0.76)$	$(-1.2,\ 1.3)$	$(-0.24,\ 0.25)$	$(-0.25,\ 0.25)$
$\mathrm{d_B} \times 10^3$	$(-0.25,\ 0.26)$	$(-0.70,\ 3.3)$	$(-0.10,\ 0.10)$	$(-0.17,\ 0.21)$
$\overline{\mathrm{d}} \times 10^3$	$(-3.2,\ 3.3)$	$(-3.6,\ 3.7)$	$(-1.5,\ 1.5)$	$(-1.2,\ 1.2)$
$\overline{\mathrm{d}}_\mathrm{B} \times 10^3$	$(-1.1,\ 1.1)$	$(-1.5,\ 1.5)$	$(-5.6,\ 5.6)$	$(-0.52,\ 0.53)$
$P_e = 1$	$p_T^e > 0$	$p_T^e > 100$ GeV	$p_T^e > 0$	$p_T^e > 100$ GeV
$d \times 10^3$	$(-0.89,\ 0.94)$	$(-10,\ 25)$	$(-0.4,\ 0.4)$	$(-3.3,\ 17)$
$\mathrm{d_B} \times 10^3$	$(-0.24,\ 0.26)$	$(-0.65,\ 1.4)$	$(-0.11,\ 0.12)$	$(-0.19,\ 0.49)$
$\overline{\mathrm{d}} \times 10^3$	$(-3.9,\ 3.9)$	$(-15,\ 15)$	$(-2.5,\ 2.7)$	$(-9.3,\ 8.1)$
$\overline{\mathrm{d}}_\mathrm{B} \times 10^3$	$(-0.97,\ 0.96)$	$(-0.97,\ 0.97)$	$(-4.5,\ 4.7)$	$(-0.31,\ 0.31)$
$P_e = -1$	$p_T^e > 0$	$p_T^e > 100$ GeV	$p_T^e > 0$	$p_T^e > 100$ GeV
$\mathrm{d} \times 10^3$	$(-0.63,\ 0.66)$	$(-0.83,\ 0.83)$	$(-0.18,\ 0.18)$	$(-0.16,\ 0.17)$
$\mathrm{d_B} \times 10^3$	$(-0.25,\ 0.26)$	$(-0.67,\ 0.66)$	$(-0.093,\ 0.094)$	$(-0.14,\ 0.14)$
$\overline{\mathrm{d}} \times 10^3$	$(-3.1,\ 3.2)$	$(-2.9,\ 3.1)$	$(-1.2,\ 1.2)$	$(-0.96,\ 0.98)$
$\overline{\mathrm{d}}_\mathrm{B} \times 10^3$	$(-1.2,\ 1.2)$	$(-2.3,\ 2.4)$	$(-0.69,\ 0.71)$	$(-0.88,\ 0.90)$

Table 1: Bounds on the CP-even d, $\mathrm{d_B}$ and CP-odd $\overline{\mathrm{d}}$, $\overline{\mathrm{d}}_\mathrm{B}$ anomalous couplings, for $\sqrt{S} = 500$ GeV and $\sqrt{S} = 1500$ GeV, e-beam polarizations $P_e = 0, 1, -1$, and for $p_T^e > 0$ and $p_T^e > 100$ GeV. In the case of $\sqrt{S} = 500$ GeV and $p_T^e > 100$ GeV, a cut on the final electron angle $\theta_e < 90^o$ is applied.

In Table 1, we report the main results for the bounds on the anomalous couplings. The bounds are obtained for $m_H = 120$ GeV, at $\sqrt{s} = 500$ GeV (with $\mathcal{L}_{int}$ =100 fb^{-1}) and $\sqrt{s} = 1500$ GeV (with $\mathcal{L}_{int}$ =1000 fb^{-1}), and for different electron-beam polarizations, $P_e = 0, 1, -1$. We assume that for each bound the only contribution to the deviation in the signal is given by the corresponding anomalous coupling, switching-off the other three anomalous contributions. A more complete analysis for the bounds on the d, $\mathrm{d_B}$, $\overline{\mathrm{d}}$,$\overline{\mathrm{d}}_\mathrm{B}$ and $\mathrm{d}_{\gamma\gamma}$, $\mathrm{d}_{\gamma Z}$, $\overline{\mathrm{d}}_{\gamma\gamma}$,$\overline{\mathrm{d}}_{\gamma Z}$ couplings can be found in [13], where the correlations between pairs of different anomalous contributions have been studied, too.

From the results of Table 1, we draw the following conclusions:

- The strongest bounds on the CP-even couplings at $\sqrt{s} = 500$ GeV are at the level of $|\mathrm{d}| \lesssim 6 \times 10^{-4}$, obtained at $P_e = -1$, and $|\mathrm{d_B}| \lesssim 2.5 \times 10^{-4}$ (with no cut on p_T^e), not depending on the e polarization. At $\sqrt{s} = 1500$ GeV, one has $|\mathrm{d}| \lesssim 1.7 \times 10^{-4}$, obtained at $P_e = -1$ and $p_T^e > 100$ GeV, and $|\mathrm{d_B}| \lesssim 1.0 \times 10^{-4}$ (with no cut on p_T^e), not depending on the e polarization.

- The strongest bounds on the CP-odd couplings at $\sqrt{s} = 500$ GeV are $|\overline{\mathrm{d}}| \lesssim 3 \times 10^{-3}$ and $|\overline{\mathrm{d}}_\mathrm{B}| \lesssim 1.0 \times 10^{-3}$, that are obtained for $P_e = -1$ and $P_e = 1$, respectively.

These bounds are quite insensitive to the cuts on p_T^e. At $\sqrt{s} = 1500$ GeV, one has $|\overline{\mathrm{d}}| \lesssim 1.0 \times 10^{-3}$ for $P_e = -1$, with $p_T^e > 100$ GeV, and $|\overline{\mathrm{d}}_{\mathrm{B}}| \lesssim 3 \times 10^{-4}$, for $P_e = 1$, with $p_T^e > 100$ GeV.

It is interesting to compare these results with other bounds obtained in the literature from different processes at future linear colliders. In particular, the processes $e^+e^- \to HZ$ and $\gamma\gamma \to H$ have been studied for a e^+e^- collider at $\sqrt{s} = 1$ TeV and with 80 fb^{-1} by Gounaris et al.. From $e^+e^- \to HZ$, they get $|\mathrm{d}| \lesssim 5 \times 10^{-3}$, $|\mathrm{d}_{\mathrm{B}}| \lesssim 2.5 \times 10^{-3}$, $|\overline{\mathrm{d}}| \lesssim 5 \times 10^{-3}$ and $|\overline{\mathrm{d}}_{\mathrm{B}}| \lesssim 2.5 \times 10^{-3}$ [18]. The process $\gamma\gamma \to H$ can do a bit better and reach the values $|\mathrm{d}| \lesssim 1 \times 10^{-3}$, $|\mathrm{d}_{\mathrm{B}}| \lesssim 3 \times 10^{-4}$, $|\overline{\mathrm{d}}| \lesssim 4 \times 10^{-3}$ and $|\overline{\mathrm{d}}_{\mathrm{B}}| \lesssim 1.3 \times 10^{-3}$, assuming a particular photon energy spectrum [19]. These analysis assume a precision of the measured production rate equal to $1/\sqrt{N}$ (with N the total number of events), and neglect possible backgrounds. In order to set the comparative potential of our process with respect to these two processes in bounding the parameters d, d_{B}, $\overline{\mathrm{d}}$, $\overline{\mathrm{d}}_{\mathrm{B}}$, we assumed $\sqrt{s} = 0.9$ TeV and (conservatively) a luminosity of 25 fb^{-1} in $e\gamma \to eH$. We then neglected any background, and assumed a precision equal to $1/\sqrt{N}$. In the case $P_e = 0$ and $p_T^e > 0$, we get $|\mathrm{d}| \lesssim 5 \times 10^{-4}$, $|\mathrm{d}_{\mathrm{B}}| \lesssim 2 \times 10^{-4}$, $|\overline{\mathrm{d}}| \lesssim 2 \times 10^{-3}$ and $|\overline{\mathrm{d}}_{\mathrm{B}}| \lesssim 8 \times 10^{-4}$.

This analysis confirms the excellent potential of the process $e\gamma \to eH$.

Following the conventions of reference [20], one can convert our constrains into upper limits for the *new physics* scale Λ that can be explored through $e\gamma \to eH$ with $\sqrt{s} \simeq$ 1.5 TeV and 10^3 fb^{-1}:

$$
\begin{aligned}
|\mathrm{d}| \lesssim 1.7 \times 10^{-4} \quad &\to \quad |\tfrac{f_{WW}}{\Lambda^2}| \lesssim 0.026 \;\; \mathrm{TeV}^{-2} \\
|\mathrm{d}_{\mathrm{B}}| \lesssim 1.0 \times 10^{-4} \quad &\to \quad |\tfrac{f_{BB}}{\Lambda^2}| \lesssim 0.015 \;\; \mathrm{TeV}^{-2} \\
|\overline{\mathrm{d}}| \lesssim 1.0 \times 10^{-3} \quad &\to \quad |\tfrac{\bar{f}_{WW}}{\Lambda^2}| \lesssim 0.15 \;\; \mathrm{TeV}^{-2} \\
|\overline{\mathrm{d}}_{\mathrm{B}}| \lesssim 3.0 \times 10^{-4} \quad &\to \quad |\tfrac{\bar{f}_{WW}}{\Lambda^2}| \lesssim 0.046 \;\; \mathrm{TeV}^{-2}
\end{aligned}
\tag{8}
$$

For $f_i \sim 1$ one can explore energy scales up to about 6, 8, 2.6 and about 4.5 TeV, respectively. At $\sqrt{s} \simeq 500$ GeV and 10^2 fb^{-1}, the corresponding constraints on the couplings are a factor 2 or 3 weaker than above (reflecting into energy scales Λ lower by a factor 1.4 or 1.7, respectively), mainly because of the smaller integrated luminosity assumed.

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